

**Scheme for fault-tolerant holonomic computation on stabilizer codes**Ognyan Oreshkov,<sup>1,2,3</sup> Todd A. Brun,<sup>2,3,4</sup> and Daniel A. Lidar<sup>2,3,4,5</sup><sup>1</sup>*Grup de Física Teòrica, Universitat Autònoma de Barcelona, Bellaterra, Barcelona 08193, Spain*<sup>2</sup>*Department of Physics, University of Southern California, Los Angeles, California 90089, USA*<sup>3</sup>*Center for Quantum Information Science & Technology, University of Southern California, Los Angeles, California 90089, USA*<sup>4</sup>*Communication Science Institute, University of Southern California, Los Angeles, California 90089, USA*<sup>5</sup>*Department of Chemistry, University of Southern California, Los Angeles, California 90089, USA*

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This paper generalizes and expands upon the work [O. Oreshkov, T. A. Brun, and D. A. Lidar, Phys. Rev. Lett. **102**, 070502 (2009)] where we introduced a scheme for fault-tolerant holonomic quantum computation (HQC) on stabilizer codes. HQC is an all-geometric strategy based on non-Abelian adiabatic holonomies, which is known to be robust against various types of errors in the control parameters. The scheme we present shows that HQC is a scalable method of computation and opens the possibility for combining the benefits of error correction with the inherent resilience of the holonomic approach. We show that with the Bacon-Shor code the scheme can be implemented using Hamiltonian operators of weights 2 and 3.

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**I. INTRODUCTION**

There are two main sources of errors in quantum computers—environment-induced decoherence and imperfect control. It has been shown that if the errors of each type are sufficiently uncorrelated and their rates are below a certain threshold, it is possible to implement reliably an arbitrarily long computational task with a modest resource overhead [1–8]. This result, known as the quantum accuracy threshold theorem, is based on the idea of quantum error correction (QEC) [9,10]—a universal software strategy to combat noise in quantum computers.

In addition to the software approach, there have also been proposals to deal with the effects of noise by hardware methods that provide robustness through their inherent properties. One such method is holonomic quantum computation (HQC) [11,12]—an adiabatic all-geometric method of computation that uses non-Abelian generalizations [13] of the Berry phase [14]. It has been shown that due to its geometric nature, this approach is robust against various types of errors in the control parameters driving the evolution [15–19] and thus provides a degree of built-in resilience at the hardware level.

In Ref. [20] HQC was combined with the method of decoherence-free subspaces (DFSs) [21–23], which was the first step toward systematic error protection in conjunction with the holonomic approach. DFSs provide *passive* protection, in the sense that no syndrome measurement and feedback are required. The DFS approach is applicable provided there is a symmetry in the system-bath coupling, which results in a sufficiently large decoupled subspace in which quantum information can be stored. Collective decoherence is a well-known example of such a symmetry (for a review see Ref. [24]). In the absence of symmetries, the addition of syndrome measurements and feedback is useful and leads to the approach based on *active* error correction [9,10]. Active error correction is also the basis of quantum fault tolerance, which is essential for the scalability of any method of computation. Even if we assume that the system is perfectly protected from environment-induced errors, when the size of the

circuit increases, errors due to imperfect operations would accumulate detrimentally unless they were corrected. Therefore, scalability of HQC requires combining the holonomic approach with active error correction.

In Ref. [25], we reported a scheme that combines HQC with the techniques for fault-tolerant computation on stabilizer codes. This demonstrated that HQC is a scalable method of computation. The scheme uses Hamiltonians that are elements of the stabilizer—or in the case of subsystem codes, elements of the stabilizer and gauge groups. Encoded gates are implemented by slowly varying the Hamiltonians along suitable paths in parameter space, such that from the point of view of the basis of the full Hilbert space, the states in each eigenspace undergo the same transversal transformation. On certain codes such as the nine-qubit Shor code [9] or its subsystem generalizations [26,27], universal computation according to our scheme can be implemented with Hamiltonians of weights 2 and 3.

This paper generalizes and expands upon the work [25]. We provide details on various proofs sketched in Ref. [25], clarify many points, and analyze properties of the scheme that were not discussed there. We examine in detail the construction for the Bacon-Shor code [26,27], discuss the adiabatic approximation for different parametrizations of the Hamiltonians, and provide explicit calculations of the holonomy in the implementation of the  $Z$  (phase-flip) gate for two different types of interpolation.

**II. PRELIMINARIES****A. Holonomic quantum computation**

Let  $\{H_\lambda\}$  be a family of Hamiltonians on an  $N$ -dimensional Hilbert space, which is continuously parametrized by a point  $\lambda$  in a control-parameter manifold  $\mathcal{M}$ . Assume that the family has the same degeneracy structure, i.e., there are no level crossings. The Hamiltonians can then be written as  $H_\lambda = \sum_{n=1}^R \varepsilon_n(\lambda) \Pi_n(\lambda)$ , where  $\{\varepsilon_n(\lambda)\}_{n=1}^R$  are the  $R$  different  $d_n$ -fold degenerate eigenvalues of  $H_\lambda$  ( $\sum_{n=1}^R d_n$

$=N$ ) and  $\Pi_n(\lambda)$  are the projectors on the corresponding eigenspaces. If the parameter  $\lambda$  is changed adiabatically, a state that initially belongs to an eigenspace of the Hamiltonian will remain in the corresponding eigenspace as the Hamiltonian evolves. The unitary evolution that results from the action of the Hamiltonian  $H(t) := H_{\lambda(t)}$  is

$$U(t) = T \exp \left[ -i \int_0^t d\tau H(\tau) \right] = \bigoplus_{n=1}^R e^{i\omega_n(t)} U_{A_n}^\lambda(t), \quad (1)$$

where  $\omega_n(t) = -\int_0^t d\tau \varepsilon_n(\lambda(\tau))$  is a dynamical phase and the intrinsically geometric operators  $U_{A_n}^\lambda(t)$  are given by the following path-ordered exponent:

$$U_{A_n}^\lambda(t) = \mathcal{P} \exp \left( \int_{\lambda(0)}^{\lambda(t)} A_n \right). \quad (2)$$

Here  $A_n$  is the Wilczek-Zee connection [13] for the  $n$ th eigenspace;  $A_n = \sum_{\mu} A_{n,\mu} d\lambda^\mu$ , where  $A_{n,\mu}$  has matrix elements [13]

$$(A_{n,\mu})_{\alpha\beta} = \langle n\alpha; \lambda | \frac{\partial}{\partial \lambda^\mu} | n\beta; \lambda \rangle. \quad (3)$$

The parameters  $\lambda^\mu$  are local coordinates on  $\mathcal{M}$  ( $1 \leq \mu \leq \dim \mathcal{M}$ ) and  $\{|n\alpha; \lambda\rangle\}_{\alpha=1}^{d_n}$  is an orthonormal basis of the  $n$ th eigenspace of the Hamiltonian at the point  $\lambda$ .

When the path  $\lambda(t)$  forms a loop  $\gamma(t)$ ,  $\gamma(0) = \gamma(T) = \lambda_0$ , the unitary matrix with components  $(U_n^\gamma)_{\alpha\beta}$  appearing in

$$\sum_{\alpha\beta} (U_n^\gamma)_{\alpha\beta} |n\alpha; 0\rangle \langle n\beta; 0| \equiv U_{A_n}^\lambda(T) = \mathcal{P} \exp \left( \oint_\gamma A_n \right) \quad (4)$$

is called the holonomy associated with the loop. In the case when the  $n$ th energy level is nondegenerate ( $d_n = 1$ ), the corresponding holonomy reduces to the Berry phase [14]. The holonomy  $U_n^\gamma$  is a geometric object which is invariant under gauge transformations corresponding to changing the basis of the  $n$ th eigenspace along the curve  $\gamma$ . The set  $\text{Hol}(A_n) = \{U_n^\gamma | \gamma \in L_{\lambda_0}(\mathcal{M})\}$ , where  $L_{\lambda_0}(\mathcal{M}) = \{\gamma: [0, T] \rightarrow \mathcal{M} | \gamma(0) = \gamma(T) = \lambda_0\}$  is the space of all loops based on  $\lambda_0$ , is a subgroup of  $U(d_n)$  called the holonomy group.

In Refs. [11, 12] it was shown that if the dimension of the control manifold is sufficiently large, quantum holonomies can be used as a means of universal quantum computation. In this approach, logical states are encoded in the degenerate eigenspace of a Hamiltonian and gates are implemented by adiabatically varying the Hamiltonian along suitable loops in the parameter manifold (for a construction of a universal set of gates, see also Ref. [28]).

We point out that many assumptions behind this simple model of HQC can be relaxed. For example, if we are interested in performing computation in the  $n$ th eigenspace of the Hamiltonian, it is sufficient that this eigenspace is adiabatically decoupled from the rest of the Hilbert space, and it is not necessary that there are no crossings between other energy levels that are separated from the  $n$ th level by energy gaps. Furthermore, in order to obtain a gauge-invariant expression for the geometric transformation taking place inside the  $n$ th eigenspace, it is not necessary that the entire Hamiltonian undergoes a cyclic change—it is enough to take the

$n$ th eigenspace around a loop. In fact, the form of the restriction of the Hamiltonian on the orthogonal complement of that subspace is irrelevant since what is important for the geometric transformation taking place inside an adiabatically decoupled eigenspace is how this subspace changes inside the full Hilbert space. More precisely, adiabatic quantum holonomies inside the  $n$ th eigenspace can be equivalently understood as arising from parallel transport of vectors along the tautological bundle whose base is the Grassmannian parametrizing the set of  $d_n$ -dimensional subspaces of the full Hilbert space rather than from parallel transport along the corresponding bundle over the full space of control parameters. We note that even the requirement for a closed loop in the Grassmannian can be relaxed using the notion of open-path holonomies [29]. The approach that we pursue in this paper can be best understood as based on closed loops in the Grassmannian, even though—with a small modification—it can be made to be exactly of the type discussed in the above formulation, where the Hamiltonian family has a fixed degeneracy structure, and gates are implemented by loops in the control manifold (see Sec. II C).

### B. Stabilizer codes and fault-tolerant computation

A large class of quantum error-correcting codes can be described by the so called stabilizer formalism [30–32]. A stabilizer  $S$  is an Abelian subgroup of the Pauli group  $\mathcal{G}_n$  on  $n$  qubits that does not contain the element  $-I$  [33]. The Pauli group consists of all possible  $n$ -fold tensor products of the Pauli matrices  $\sigma^x \equiv X$ ,  $\sigma^y \equiv Y$ ,  $\sigma^z \equiv Z$  together with the multiplicative factors  $\pm 1$ ,  $\pm i$ . The stabilizer code corresponding to  $S$  is the subspace of all states  $|\psi\rangle$  that are left invariant under the action of every operator in  $S$  ( $G|\psi\rangle = |\psi\rangle$ ,  $\forall G \in S$ ). It is easy to see that the stabilizer of a code encoding  $k$  qubits into  $n$  has  $n-k$  generators. Recently, a more general notion of codes has been introduced—*subsystem* or *operator* codes [34, 35]—that employs the most general encoding of information, encoding in subsystems [36–38]. In the case of a single subsystem, the Hilbert space decomposes as  $\mathcal{H} = \mathcal{H}^A \otimes \mathcal{H}^B \oplus \mathcal{K}$ , where  $\mathcal{H}^A$  is the subsystem in which the logical information is encoded,  $\mathcal{H}^B$  is the *gauge* subsystem, and  $\mathcal{K}$  is the rest of the Hilbert space. For operator stabilizer codes, the stabilizer leaves the subspace  $\mathcal{H}^A \otimes \mathcal{H}^B$  invariant but the encoded information is invariant also under operations that act on the gauge subsystem. An operator stabilizer code encoding  $k$  qubits into  $n$  with  $r$  gauge qubits has  $n-r-k$  stabilizer generators, while the gauge group has  $2r$  generators [39]. According to the error-correction condition for stabilizer codes [33, 39], a set of errors  $\{E_i\}$  in  $\mathcal{G}_n$  (which without loss of generality are assumed to be Hermitian) is correctable by the code if and only if, for all  $i$  and  $j$ ,  $E_i E_j$  anticommutes with at least one element of  $S$  or otherwise belongs to  $S$  or to the gauge group. In this paper we will be concerned with stabilizer codes for the correction of single-qubit errors and the techniques for fault-tolerant computation [1–8] on such codes.

A quantum information processing scheme is called fault tolerant if a single error occurring during the implementation of any given operation introduces at most one error per block

of the code [7]. This property has to apply for unitary gates as well as measurements, including those that constitute the error-correcting operations themselves. Fault-tolerant schemes for computation on stabilizer codes generally depend on the code being used—some codes, such as the Bacon-Shor subsystem codes [26,27], for example, are better suited for fault-tolerant computation than others [40]. In spite of these differences, however, it has been shown that fault-tolerant information processing is possible on any stabilizer code [6,7]. The general procedure can be described briefly as follows. DiVincenzo and Shor [2] demonstrated how to perform fault-tolerant measurements of the stabilizer for any stabilizer code. Their method makes use of an approach introduced by Shor [1], which involves the “cat” state  $(|0\dots 0\rangle + |1\dots 1\rangle)/\sqrt{2}$  which can be prepared and verified with a satisfactory precision. As pointed out by Gottesman [7], by the same method one can measure any operator in the Pauli group. Since the encoded  $X$ ,  $Y$ , and  $Z$  operators belong to the Pauli group for any stabilizer code [7], one can prepare fault tolerantly various superpositions of the logical basis states  $|\bar{0}\rangle$  and  $|\bar{1}\rangle$ , such as  $|\bar{\pm}\rangle = (|\bar{0}\rangle + |\bar{1}\rangle)/\sqrt{2}$ , for example. The latter can be used to implement fault tolerantly the encoded phase and Hadamard gates as long as a fault-tolerant controlled-NOT (C-NOT) gate is available [7]. Gottesman showed how the C-NOT gate can be implemented fault tolerantly by first applying a transversal operation on four encoded qubits and then measuring the encoded  $X$  operator on two of them. Finally, for universal computation one needs a gate outside of the Clifford group, e.g., the Toffoli gate. The Toffoli construction was demonstrated first by Shor in Ref. [1] for a specific type of codes—those obtained from doubly-even self-dual classical codes by the Calderbank-Shor-Steane (CSS) construction [41,42]. Gottesman [6] showed that a transversal implementation of the same procedure exists for any stabilizer code.

### C. Overview of the scheme

Note that the described method for universal fault-tolerant computation on stabilizer codes uses almost exclusively transversal operations—these are operations for which each qubit in a block interacts only with the corresponding qubit from another block or from a special ancillary state such as Shor’s cat state (see also the methods of Steane [43] and Knill [44]). However, transversal operations are not the most general class of operations that do not lead to propagation of errors. For example, every transversal operation in a given fault-tolerant protocol can be substituted by the same operation followed by an operation that multiplies each syndrome subspace by a different phase, and the resultant protocol will still be fault tolerant. This can be easily seen from the fact that if after a transversal operation the state is correctable, then it will still be correctable after multiplying each syndrome subspace by a phase because the correction procedure involves a projection on one of the syndrome subspaces and thus the overall phase in that subspace is irrelevant. An operation that is equal to a transversal operation followed by a transformation on the gauge subsystem can be similarly seen to be fault tolerant. It is these more general fault-tolerant

transformations by means of which we will realize fault-tolerant HQC.

To explain the main idea behind our approach, let us consider the case of standard (subspace) stabilizer codes first. Our goal will be to find a holonomic realization of a universal set of encoded gates by adiabatically transporting the code space along suitable loops via sequences of elementary fault-tolerant transformations of the above type. The scheme we will present can be roughly described as follows. We choose as a starting Hamiltonian an element of the stabilizer and vary this Hamiltonian in an adiabatic manner along appropriate paths so that from the point of view of the basis of the full Hilbert space, the vectors in both eigenspaces of the Hamiltonian undergo the same transversal transformation. Under this procedure, each eigenspace will acquire a dynamical phase corresponding to the energy of that eigenspace, but since the Hamiltonian is an element of the stabilizer, these phases will amount to relative phases between different syndrome subspaces and they would be projected out if a measurement of the syndrome is performed. So from the point of view of the basis of the full Hilbert space, the overall transformation under this procedure is of the general fault-tolerant type we described. Since the code space is a subspace of an eigenspace of the Hamiltonian, it will be effectively transformed by the corresponding transversal operation.

The standard fault-tolerant procedures provide prescriptions of how to implement any encoded gate by a sequence of elementary transversal operations. Therefore, if we make the code space follow an appropriate sequence of transversal operations in the described adiabatic manner, when we complete an encoded operation we will have taken the code space around a loop whose associated holonomy is equal to the encoded operation. A simple way to see that this is indeed a holonomy is to notice that if we track the initial code space as it evolves, it undergoes a loop in the Grassmannian since at the end we complete an encoded gate. Furthermore, at all times the Hamiltonian acts trivially on the code space so that all states inside it acquire the same dynamical phase. Hence the nontrivial transformation resulting inside the code space must be geometric. Thus by following precisely the sequence of transversal operations and measurements that are used in a given dynamical fault-tolerant scheme, we obtain a scheme that implements logical gates through adiabatic holonomies and at the same time is fault tolerant.

In the case of subsystem codes, the code subsystem can be thought of as a collection of subspaces, each of which contains the same redundant information. The relative phases between each of these subspaces are gauge degrees of freedom. Applying a particular encoded gate is equivalent to applying the same gate in each subspace. In this case, our scheme can use Hamiltonians that are either elements of the stabilizer or elements of the gauge group. If the Hamiltonian is an element of the stabilizer, all “redundant” subspaces inside a given syndrome subspace belong to a single eigenspace of the Hamiltonian and all of them will undergo the transversal operation that effectively transforms that eigenspace. If the Hamiltonian is an element of the gauge group, then some of the subspaces of interest will belong to the ground space while others will belong to the excited space.

However, since the scheme implements the same transversal operation in each eigenspace, all subspaces by construction will undergo the same transversal operation. The relative dynamical phases acquired between subspaces in the ground and excited eigenspaces amount to a gauge transformation that does not affect the fault tolerance of the operation. Thus by following the same sequence of transversal operations and measurements as in a standard dynamical scheme, we obtain a fault-tolerant holonomic realization of encoded computation on subsystem codes.

From the point of view of the full Hilbert space, this method performs transformations of the generalized fault-tolerant type we described earlier. This is what ensures the fault tolerance of the method. From a geometrical point of view, this corresponds to transporting each of the different syndrome subspaces around a loop such that all of them simultaneously undergo the same geometric transformation (in the case of subsystem codes, the same applies for all redundant subspaces inside each syndrome space). The statement that all subspaces simultaneously undergo the same geometric transformation makes sense with respect to a particular choice of the instantaneous basis in each subspace. For any choice of basis in the code space at a given point along the loop, there is a preferred choice of bases inside the rest of the syndrome spaces that is determined by the notion of correctable errors. Correctable errors can be thought of as transitions from the code space to the other syndrome spaces which can be undone if error correction is applied. These errors therefore map the basis from the code space to particular bases in the error spaces such that with respect to these bases, a state that has undergone a correctable error has the same form as the form that the nonerroneous state has with respect to the basis of the code space. Our scheme applies the same geometric transformation, in this sense, in all syndrome subspaces.

Note that the present approach differs from the original HQC model [11,12] in that it computes in several subspaces at the same time. Another difference from the original model is that we do not use a single family of isodegenerate Hamiltonians. This is because for simplicity we use Hamiltonians that are equal to a single element of the stabilizer or the gauge group at a given time, and we change that element along different portions of the loop. Thus, if the Hamiltonians are elements of the stabilizer, a particular syndrome subspace may belong to the ground space of the Hamiltonian during one portion of the loop and to the excited space during another one. Therefore, the holonomies in our scheme can be most naturally understood as resulting from parallel transport along loops in the tautological bundle over the Grassmannian (this is the fiber bundle whose base is the Grassmannian and whose fibers are the subspaces corresponding to the different points in the Grassmannian) rather than loops in a bundle over a space parametrizing an isodegenerate family of Hamiltonians. In the case of subsystem codes, if along the loop we change between Hamiltonians which are noncommuting elements of the gauge group, the redundant subspaces that constitute a logical subsystem may seem to undergo dynamical transformations in addition to the geometric ones. However, these dynamical changes are equivalent to gauge transformations and do not affect the

workings of the scheme. We could modify our scheme so that it uses a Hamiltonian that separates all of the subspaces of interest by energetic gaps and adiabatically transports each of these subspaces along the same path that it would follow under the scheme we described. Then the holonomy resulting in each eigenspace could be understood as being of the original HQC type. However, this is unnecessary since the dynamical phases are irrelevant for the workings of our model.

As single-qubit unitaries together with the C-NOT gate form a universal set of gates, fault-tolerant computation can be realized entirely in terms of single-qubit operations and C-NOT operations between qubits from different blocks, assuming that the cat state can be prepared reliably. Hence our task will be to find adiabatic realizations of these operations, as well as of the operations for preparing and verifying the cat state. Then these operations can be used as building blocks to implement fault-tolerant HQC according to the idea described in this section.

### III. SCHEME

Consider a  $[[n, 1, r, 3]]$  stabilizer code. This is a code that encodes one qubit into  $n$ , has  $r$  gauge qubits, and can correct arbitrary single-qubit errors. To perform a holonomic operation on this code, we need a nontrivial starting Hamiltonian that leaves the code space or code subsystem invariant. It is easy to verify that the only Hamiltonians that satisfy this property are linear combinations of the elements of the stabilizer and, in the case of subsystem codes, elements of the gauge group.

Note that the stabilizer and the gauge group transform during the course of the computation under the operations being applied. At any stage when we complete an encoded operation, they return to their initial forms. Our scheme will follow the same transversal operations as those used in a standard dynamical fault-tolerant scheme, but as we explained in Sec. II C, in addition we will have extra dynamical phases that multiply each syndrome subspace or are equivalent to gauge transformations. However, it is easy to see that these phases do not affect the way the stabilizer or the gauge group transform, so we can omit them from our analysis of the transformation of these groups.

During the implementation of a standard encoded gate, the Pauli group  $\mathcal{G}_n$  on a given codeword may change in such a way that it acts on other codewords, but it can be verified that this “spreading” can be limited to at most four other codewords including the cat state. This is because the encoded C-NOT gate can be implemented fault tolerantly on any stabilizer code by a transversal operation on four encoded qubits [6], and any encoded Clifford gate can be realized using only the encoded C-NOT, provided that we are able to do fault-tolerant measurements (the encoded Clifford group is generated by the encoded Hadamard, phase, and C-NOT gates). Encoded gates outside of the Clifford group, such as the encoded  $\pi/8$  or Toffoli gates, can be implemented fault tolerantly using encoded C-NOT gates conditioned on the qubits in a cat state, so they may require transversal operations on a total of five blocks. More precisely, the fault-tolerant implementation of the Toffoli gate requires the preparation of

a special state of three encoded qubits [1], which involves a sequence of conditional encoded phase operations and conditional encoded C-NOT operations with conditioning on the qubits in a cat state [6]. But the encoded phase gate has a universal implementation using an encoded C-NOT between the qubit and an ancilla, so the conditional phase gate may require applying a conditional encoded C-NOT. The procedure for implementing an encoded  $\pi/8$  gate involves applying an encoded  $SX$  gate conditioned on the qubits in a cat state [45] ( $S$  denotes the phase gate), but the encoded  $S$  gate generally also involves an encoded C-NOT on the qubit and an ancilla, so it may also require the interaction of four blocks. For CSS codes, however, the spreading of the Pauli group that acts on a given block can be limited to a total of three blocks during the implementation of a basic encoded operation.

It also should be noted that fault-tolerant encoded Clifford operations can be implemented using only Clifford gates on the physical qubits [6]. These operations transform the stabilizer and the gauge group into subgroups of the Pauli group, and their elements remain in the form of tensor products of Pauli matrices. The fault-tolerant implementation of encoded gates outside of the Clifford group, however, involves operations that take these groups outside of the Pauli group. We will, therefore, consider separately two cases: encoded operations in the Clifford group and encoded operations outside of the Clifford group.

### A. Encoded operations in the Clifford group

In Ref. [6] it was shown that every encoded operation in the Clifford group can be implemented fault tolerantly using Clifford gates on physical qubits. The Clifford group is generated by the Hadamard, phase, and C-NOT gates, but in addition to these gates, we will also demonstrate the adiabatic implementation of the  $X$  and  $Z$  gates which are standard for quantum computation. We will restrict our attention to implementing single-qubit unitaries on the first qubit in a block, as well as C-NOT operations between the first qubits in two blocks. The operations on the rest of the qubits can be obtained analogously.

#### 1. Single-qubit unitary operations

In order to implement a single-qubit operation on the first qubit in a block, we will choose as a starting Hamiltonian an element of the stabilizer (with a minus sign) or an element of the gauge group that acts nontrivially on that qubit. Since we are considering codes that can correct arbitrary single-qubit errors, one can always find an element  $\hat{G}$  of the initial stabilizer or the initial gauge group that has a factor  $\sigma^0=I$ ,  $\sigma^1=X$ ,  $\sigma^2=Y$ , or  $\sigma^3=Z$  acting on the first qubit, i.e.,

$$\hat{G} = \sigma^i \otimes \tilde{G}, \quad i = 0, 1, 2, 3, \quad (5)$$

where  $\tilde{G}$  is a tensor product of Pauli matrices and the identity on the remaining  $n-1$  qubits. It can be verified that under Clifford gates the stabilizer and the gauge group transform in such a way that this is always the case except that the factor  $\tilde{G}$  may spread to qubits in other blocks. From now on, we will use “hat” to denote operators on all these qubits and

“tilde” to denote operators on all the qubits except the first one.

Without loss of generality we will assume that the chosen stabilizer or gauge-group element for that qubit has the form

$$\hat{G} = Z \otimes \tilde{G}. \quad (6)$$

As initial Hamiltonian, we will take the operator

$$\hat{H}(0) = -\hat{G} = -Z \otimes \tilde{G}. \quad (7)$$

Thus, if  $\hat{G}$  is an element of the stabilizer, the code space will belong to the ground space of  $\hat{H}(0)$ . Our goal is to find paths in the space of parameters of the Hamiltonian such that when the Hamiltonian is varied adiabatically along these paths, each of its eigenspaces undergoes a transformation which is equivalent to that caused by a single-qubit operation on the first qubit.

*Proposition 1.* If the initial Hamiltonian (7) is varied adiabatically so that only the factor acting on the first qubit changes,

$$\hat{H}(t) = -H(t) \otimes \tilde{G}, \quad (8)$$

where

$$\text{Tr}\{H(t)\} = 0, \quad (9)$$

the transformation that each of the eigenspaces of this Hamiltonian undergoes will be equivalent to that resulting from a local unitary on the first qubit up to a global phase, i.e., the geometric part  $\hat{U}_g(t) = U_{A_0}^\lambda \oplus U_{A_1}^\lambda$  of the overall unitary, where  $U_{A_i}^\lambda$ ,  $i=0, 1$ , are the path-ordered exponents [Eq. (2)] corresponding to the ground and excited spaces, respectively, will be equal to a local unitary on the first qubit,  $\hat{U}_g(t) = U(t) \otimes \tilde{I}$ .

*Proof.* Observe that Eq. (8) can be written as

$$\hat{H}(t) = H(t) \otimes \tilde{P}_0 - H(t) \otimes \tilde{P}_1, \quad (10)$$

where

$$\tilde{P}_0 = \frac{\tilde{I} - \tilde{G}}{2}, \quad \tilde{P}_1 = \frac{\tilde{I} + \tilde{G}}{2} \quad (11)$$

are orthogonal complementary projectors. The evolution driven by  $\hat{H}(t)$  is therefore

$$\hat{U}(t) = U_0(t) \otimes \tilde{P}_0 + U_1(t) \otimes \tilde{P}_1, \quad (12)$$

where

$$U_0(t) = \mathcal{T} \exp \left[ -i \int_0^t H(\tau) d\tau \right],$$

$$U_1(t) = \mathcal{T} \exp \left[ i \int_0^t H(\tau) d\tau \right]. \quad (13)$$

Let  $|\phi_0(t)\rangle$  and  $|\phi_1(t)\rangle$  be the instantaneous ground and excited states of  $H(t)$  with eigenvalues  $E_0(t) = -E(t)$ ,  $E_1(t) = E(t)$  [ $E(t) > 0$ ]. Using Eq. (1) for the expression in (13), we obtain that in the adiabatic limit,

$$U_j(t) = e^{i\omega(t)}U_{A_j}(t) \oplus e^{-i\omega(t)}U_{A_j}(t), \quad j = 0, 1, \quad (14)$$

where  $\omega(t) = \int_0^t d\tau E(\tau)$  and

$$U_{A_j}(t) = \exp \left[ \int_0^t d\tau \langle \phi_j(\tau) | \frac{d}{d\tau} | \phi_j(\tau) \rangle \right] | \phi_j(t) \rangle \langle \phi_j(0) |, \quad j = 0, 1. \quad (15)$$

The projectors on the ground and excited eigenspaces of  $\hat{H}(0)$  are

$$\hat{P}_0 = | \phi_0(0) \rangle \langle \phi_0(0) | \otimes \tilde{P}_0 + | \phi_1(0) \rangle \langle \phi_1(0) | \otimes \tilde{P}_1 \quad (16)$$

and

$$\hat{P}_1 = | \phi_1(0) \rangle \langle \phi_1(0) | \otimes \tilde{P}_0 + | \phi_0(0) \rangle \langle \phi_0(0) | \otimes \tilde{P}_1, \quad (17)$$

respectively. Using Eqs. (14) and (15), one can see that the effect of the unitary in (12) on each of these projectors is

$$\hat{U}(t)\hat{P}_0 = e^{i\omega(t)}[U_{A_0}(t) \oplus U_{A_1}(t)] \otimes \tilde{I}\hat{P}_0, \quad (18)$$

$$\hat{U}(t)\hat{P}_1 = e^{-i\omega(t)}[U_{A_0}(t) \oplus U_{A_1}(t)] \otimes \tilde{I}\hat{P}_1, \quad (19)$$

i.e., up to an overall dynamical phase its effect on each of the eigenspaces is the same as that of the unitary

$$\hat{U}_g(t) = U(t) \otimes \tilde{I}, \quad (20)$$

where

$$U(t) = U_{A_0}(t) \oplus U_{A_1}(t). \quad (21)$$

This completes the proof.

We next show how by suitably choosing  $H(t)$  we can implement all necessary single-qubit gates. We will identify a set of points in parameter space such that by interpolating between these points we can draw various paths resulting in the desired transformations. We remark that if a path does not form a loop, the geometric transformation in (21) could be associated with an open-path holonomy taking place inside each eigenspace of the Hamiltonian, provided that the final eigenspaces have nonzero overlap with the corresponding initial ones [29] (see Appendix B).

Consider the single-qubit unitary operator

$$V^{\theta\pm} = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & \mp e^{-i\theta} \\ \pm e^{i\theta} & 1 \end{pmatrix}, \quad (22)$$

where  $\theta$  is a real parameter. Note that  $V^{\theta\mp} = (V^{\theta\pm})^\dagger$ . Define the following single-qubit Hamiltonian:

$$H^{\theta\pm} \equiv V^{\theta\pm} Z V^{\theta\mp}. \quad (23)$$

Let  $H(t)$  in Eq. (8) be a Hamiltonian which interpolates between  $H(0)=Z$  and  $H(T)=H^{\theta\pm}$  (up to a factor) as follows:

$$H(t) = f(t)Z + g(t)H^{\theta\pm} \equiv H_{f,g}^{\theta\pm}(t), \quad (24)$$

where  $f(0) > 0$ ,  $g(T) > 0$ , and  $f(T) = g(0) = 0$ . To simplify our notation, we will drop the indices  $f$  and  $g$  of the Hamiltonian since the exact form of these functions is not important for our analysis as long as they are sufficiently smooth (see dis-

cussion below). This Hamiltonian has eigenvalues  $\pm \sqrt{f(t)^2 + g(t)^2}$  and its energy gap is nonzero unless the entire Hamiltonian vanishes.

*Proposition 2.* In the adiabatic limit, the Hamiltonian in (8) with  $H(t) = H^{\theta\pm}(t)$  gives rise to the geometric transformation  $V^{\theta\pm} \otimes \tilde{I}$ .

The proof of this proposition is presented in Appendix A.

We will use this result to construct a set of standard gates by sequences of operations of the form  $V^{\theta\pm}$ , which can be generated by interpolations of the type given in Eq. (24) run forward or backward. For single-qubit gates in the Clifford group, we will only need three values of  $\theta$ : 0,  $\pi/2$ , and  $\pi/4$ . For completeness, however, we will also demonstrate how to implement the  $\pi/8$  gate, which together with the Hadamard gate is sufficient to generate any single-qubit unitary transformation [45]. For this we will need  $\theta = \pi/8$ . Note that

$$H^{\theta\pm} = \pm (\cos \theta X + \sin \theta Y), \quad (25)$$

so for these values of  $\theta$  we have  $H^{0\pm} = \pm X$ ,  $H^{\pi/2\pm} = \pm Y$ ,  $H^{\pi/4\pm} = \pm (\frac{1}{\sqrt{2}}X + \frac{1}{\sqrt{2}}Y)$ , and  $H^{\pi/8\pm} = \pm (\cos \frac{\pi}{8}X + \sin \frac{\pi}{8}Y)$ .

Consider the following adiabatic interpolations:

$$-Z \otimes \tilde{G} \rightarrow -Y \otimes \tilde{G} \rightarrow Z \otimes \tilde{G}. \quad (26)$$

According to the above result, the first interpolation yields the transformation  $V^{\pi/2+}$ . The second interpolation can be regarded as the inverse of  $Z \otimes \tilde{G} \rightarrow -Y \otimes \tilde{G}$  which is equivalent to  $-Z \otimes \tilde{G} \rightarrow Y \otimes \tilde{G}$  since  $\hat{H}(t)$  and  $-\hat{H}(t)$  yield the same geometric transformations. Thus the second interpolation results in  $(V^{\pi/2-})^\dagger = V^{\pi/2+}$ . The net result is therefore  $V^{\pi/2+}V^{\pi/2+} = iX$ . We see that up to a global phase, the geometric part of the transformation resulting from the above sequence is equal to the single-qubit  $X$  gate.

Similarly, one can verify that the  $Z$  gate can be realized via the loop

$$-Z \otimes \tilde{G} \rightarrow -X \otimes \tilde{G} \rightarrow Z \otimes \tilde{G} \rightarrow Y \otimes \tilde{G} \rightarrow -Z \otimes \tilde{G}. \quad (27)$$

The phase gate can be realized by applying

$$-Z \otimes \tilde{G} \rightarrow -\left(\frac{1}{\sqrt{2}}X + \frac{1}{\sqrt{2}}Y\right) \otimes \tilde{G} \rightarrow Z \otimes \tilde{G}, \quad (28)$$

followed by the  $X$  gate.

The Hadamard gate can be realized by first applying  $Z$ , followed by

$$-Z \otimes \tilde{G} \rightarrow -X \otimes \tilde{G}. \quad (29)$$

Finally, the  $\pi/8$  gate can be implemented by first applying  $Y = iXZ$ , followed by

$$Z \otimes \tilde{G} \rightarrow -\left(\cos \frac{\pi}{8}X + \sin \frac{\pi}{8}Y\right) \otimes \tilde{G} \rightarrow -Z \otimes \tilde{G}. \quad (30)$$

*Comment.* We emphasize that the elementary gates we construct here are not holonomies associated with the code space or the error subspaces. The holonomy in a given subspace is a transformation whose components are defined with respect to a basis of that subspace, while these elementary

gates are the geometric parts of the unitaries resulting in the full Hilbert space under the described adiabatic evolutions. As explained in Sec. II C, only after we compose a suitable sequence of such elementary gates do we obtain a closed-loop holonomy in the code space (or the code subsystem). The non-Abelian holonomies through which we perform HQC in the code space and the error subspaces are the universal set of encoded gates that we would obtain by composing the set of elementary one- and two-qubit gates (combined with measurements if necessary) according to the rules of a given dynamical fault-tolerant scheme.

## 2. Note on the adiabatic condition

Before we show how to implement the C-NOT gate between two qubits, let us comment on the conditions under which the adiabatic approximation assumed in the above operations is satisfied. Because of the form [Eq. (12)] of the overall unitary, the adiabatic approximation depends on the extent to which each of the unitaries in Eq. (13) approximates expression (14). The latter depends only on the adiabatic properties of the nondegenerate two-level Hamiltonian  $H(t)$ . For such a Hamiltonian, the simple version of the adiabatic condition [46] reads

$$\frac{\epsilon}{\Delta^2} \ll 1, \quad (31)$$

where

$$\epsilon = \max_{0 \leq t \leq T} \left| \langle \phi_1(t) | \frac{dH(t)}{dt} | \phi_0(t) \rangle \right|, \quad (32)$$

and

$$\Delta = \min_{0 \leq t \leq T} [E_1(t) - E_0(t)] = \min_{0 \leq t \leq T} 2E(t) \quad (33)$$

is the minimum energy gap of  $H(t)$ .

Along the segments of the parameter paths we described, the Hamiltonian is of the form given in Eq. (24) and its derivative is

$$\frac{dH^{\theta \pm}(t)}{dt} = \frac{df(t)}{dt} Z + \frac{dg(t)}{dt} H^{\theta \pm}, \quad 0 < t < T. \quad (34)$$

This derivative is well defined as long as  $\frac{df(t)}{dt}$  and  $\frac{dg(t)}{dt}$  are well defined. The curves we described, however, may not be differentiable at the points connecting two segments. In order for the Hamiltonian in Eq. (24) that interpolates between these points to be differentiable, the functions  $f(t)$  and  $g(t)$  have to satisfy  $\frac{df(T)}{dt} = 0$  and  $\frac{dg(0)}{dt} = 0$ . This means that the change in the Hamiltonian slows down to zero at the end of each segment (except for a possible change in its strength) and increases again from zero along the next segment. We point out that when the Hamiltonian stops changing, we can turn it off completely by decreasing its strength. This can be done arbitrarily fast and it would not affect a state that belongs to an eigenspace of the Hamiltonian. Similarly, we can turn on another Hamiltonian for the implementation of a different operation.

The above condition guarantees that the adiabatic approximation is satisfied with precision  $O[(\frac{\epsilon}{\Delta^2})^2]$ . It is known,

however, that under certain conditions on the Hamiltonian, we can obtain better results [47,48]. Let us write the Schrödinger equation as

$$i \frac{d}{dt} |\psi(t)\rangle = H(t) |\psi(t)\rangle \equiv \frac{1}{\epsilon} \bar{H}(t) |\psi(t)\rangle, \quad (35)$$

where  $\epsilon > 0$  is small. Assume that  $\bar{H}(t)$  is smooth and all its derivatives vanish at the end points  $t=0$  and  $t=T$  (note that this means that  $\bar{H}(t)$  is nonanalytic at these points unless it is constant; an alternative strategy is to consider analytic Hamiltonians with a finite number of vanishing derivatives at  $t=0, T$  [48]). Then if we keep  $\bar{H}(t)$  fixed and vary  $\epsilon$ , the adiabatic error would scale superpolynomially with  $\epsilon$ , i.e., the error will decrease with  $\epsilon$  faster than  $O(\epsilon^N)$  for any  $N$  [47]. [Notice that  $\frac{\epsilon}{\Delta^2} \propto \epsilon$ , i.e., the error according to the standard adiabatic approximation is  $O(\epsilon^2)$ .]

In our case, the smoothness condition translates directly to the functions  $f(t)$  and  $g(t)$ . For any smooth  $f(t)$  and  $g(t)$  we can further ensure that the condition at the end points is satisfied by the reparametrization  $f(t) \rightarrow f[y(t)]$ ,  $g(t) \rightarrow g[y(t)]$ , where  $y(t)$  is a smooth function of  $t$  which satisfies  $y(0)=0$ ,  $y(T)=T$  and has vanishing derivatives at  $t=0$  and  $t=T$ . Then by slowing down the change in the Hamiltonian by a constant factor  $\epsilon$ , which amounts to increasing the total time  $T$  by a factor  $1/\epsilon$ , we can decrease the error superpolynomially in  $\epsilon$ . We will use this result to obtain a low-error interpolation in Sec. IV, where we estimate the time needed to implement a geometric gate with a certain precision.

## 3. C-NOT gate

The stabilizer (or gauge group) on multiple blocks of the code is a direct product of the stabilizers (or gauge groups) of the individual blocks. Therefore, from Eq. (5) it follows that one can always find an element of the initial stabilizer or gauge group on multiple blocks that has any desired combination of factors  $\sigma^i$ ,  $i=0, 1, 2, 3$ , on the first qubits in these blocks. It can be verified that applying transversal Clifford operations on the blocks does not change this property. Therefore, we can find an element of the stabilizer or the gauge group that has the form given in Eq. (6), where the factor  $Z$  acts on the target qubit and  $\tilde{G}$  acts trivially on the control qubit. We now explain how to implement the C-NOT gate geometrically starting from such a Hamiltonian.

Notice that a Hamiltonian of the form

$$\hat{H}(t) = |0\rangle\langle 0|^c \otimes H_0(t) \otimes \tilde{G} + |1\rangle\langle 1|^c \otimes H_1(t) \otimes \tilde{G}, \quad (36)$$

where the superscript  $c$  denotes the control qubit, gives rise to the unitary transformation

$$\hat{U}(t) = |0\rangle\langle 0|^c \otimes \hat{U}_0(t) + |1\rangle\langle 1|^c \otimes \hat{U}_1(t), \quad (37)$$

where

$$\hat{U}_{0,1}(t) = \mathcal{T} \exp \left[ -i \int_0^t d\tau H_{0,1}(\tau) \otimes \tilde{G} \right]. \quad (38)$$

If  $H_0(t)$  and  $H_1(t)$  have the same nondegenerate instantaneous spectra and  $\text{Tr}\{H_{0,1}(t)\}=0$ , then from Eq. (1) and proposition 1 it follows that in the adiabatic limit each of the eigenspaces of  $\hat{H}(t)$  will undergo the geometric transformation,

$$\hat{U}_g(t) = |0\rangle\langle 0|^c \otimes V_0(t) \otimes \tilde{I} + |1\rangle\langle 1|^c \otimes V_1(t) \otimes \tilde{I}, \quad (39)$$

where  $V_{0,1}(t) \otimes \tilde{I}$  are the geometric transformations generated by  $H_{0,1}(t) \otimes \tilde{G}$  according to proposition 1.

Our goal is to find  $H_0(t)$  and  $H_1(t)$ ,  $H_0(0)=H_1(0)=Z$ , such that at the end of the transformation, the geometric unitary in Eq. (39) will be equal to the C-NOT gate. In other words, we want  $V_0(2T)=I$  and  $V_1(2T)=X$  (here we have chosen the total time of evolution to be  $2T$  for convenience).

We already saw how to generate geometrically the  $X$  gate up to a phase [Eq. (26)]. We can use the same Hamiltonian in place of  $H_1(t)$ ,

$$H_1(t) = \begin{cases} H^{\pi/2+}(t), & 0 \leq t \leq T \\ H^{\pi/2-}(2T-t), & T \leq t \leq 2T. \end{cases} \quad (40)$$

Now we want to find a Hamiltonian  $H_0(t)$  with the same spectrum as  $H_1(t)$ , which gives rise to a trivial geometric transformation,  $V_0(2T)=I$  (possibly up to a phase, which can be undone later). Since all Hamiltonians of the type  $H^{\theta\pm}(t)$  have the same instantaneous spectrum [for fixed  $f(t)$  and  $g(t)$ ], we can simply choose

$$H_0(t) = \begin{cases} H^{\pi/2+}(t), & 0 \leq t \leq T \\ H^{\pi/2+}(2T-t), & T \leq t \leq 2T, \end{cases} \quad (41)$$

which corresponds to applying a given transformation from  $t=0$  to  $t=T$  and then undoing it (running it backward) from  $t=T$  to  $t=2T$ . This results exactly in  $V_0(2T)=I$ .

Since, as we saw in Sec. III A 1, the Hamiltonian  $H_1(t) \otimes \tilde{G}$  gives rise to the geometric transformation  $iX \otimes \tilde{I}$ , the above choice for the Hamiltonians in Eqs. (40) and (41) in Eq. (36) will result in the geometric transformation

$$|0\rangle\langle 0|^c \otimes I \otimes \tilde{I} + i|1\rangle\langle 1|^c \otimes X \otimes \tilde{I}, \quad (42)$$

which is the desired C-NOT gate up to a phase gate on the control qubit. We can correct the phase by applying the inverse of the phase gate to the control qubit either before or after the described transformation.

Notice that from  $t=0$  to  $t=T$  the Hamiltonians in Eqs. (40) and (41) are identical, i.e., during this period the Hamiltonian in Eq. (36) has the form

$$I^c \otimes H^{\pi/2+}(t) \otimes \tilde{G}, \quad (43)$$

so we are simply applying the single-qubit operation  $V^{\pi/2+}$  to the target qubit according to the method for single-qubit gates described before. It is straightforward to verify that during the second period, from  $t=T$  to  $t=2T$ , the Hamiltonian in Eq. (36) realizes the interpolation

$$-I^c \otimes Y \otimes \tilde{G} \rightarrow -Z^c \otimes Z \otimes \tilde{G}, \quad (44)$$

which is understood as in Eq. (24).

To summarize, the C-NOT gate can be implemented by first applying the inverse of the phase gate ( $S^\dagger$ ) on the control qubit, as well as the transformation  $V^{\pi/2+}$  on the target qubit, followed by transformation (44). Due to the form of  $\hat{H}(t)$  in Eq. (36), the extent to which the adiabatic approximation is satisfied during this transformation depends only on the adiabatic properties of the single-qubit Hamiltonians  $H^{\pi/2\pm}(t)$  which we discussed in Sec. III A 2.

Our construction allowed us to prove the resulting geometric transformations without explicitly calculating the path-ordered integrals [Eq. (2)]. It may be instructive, however, to demonstrate this calculation for at least one of the gates we described. In Appendix B, we present an explicit calculation of the geometric transformation for the  $Z$  gate for the following two cases:  $f(t)=1-\frac{t}{T}$ ,  $g(t)=\frac{t}{T}$  (linear interpolation);  $f(t)=\cos\frac{\pi t}{2T}$ ,  $g(t)=\sin\frac{\pi t}{2T}$  (unitary interpolation).

### B. Encoded operations outside of the Clifford group

For universal fault-tolerant computation we also need at least one encoded gate outside of the Clifford group. The fault-tolerant implementation of such gates is based on the preparation of a special encoded state [1,5,6,45,49] which involves a measurement of an encoded operator in the Clifford group. For example, the  $\pi/8$  gate requires the preparation of the encoded state  $\frac{|0\rangle+\exp(i\pi/4)|1\rangle}{\sqrt{2}}$ , which can be realized by measuring the encoded operator  $e^{-i\pi/4}SX$  [45]. Equivalently, the state can be obtained by applying the encoded operation  $RS^\dagger$ , where  $R$  denotes the Hadamard gate, on the encoded state  $\frac{\cos(\pi/8)|0\rangle+\sin(\pi/8)|1\rangle}{\sqrt{2}}$  which can be prepared by measuring the encoded Hadamard gate [5]. The Toffoli gate requires the preparation of the three-qubit encoded state  $\frac{|000\rangle+|010\rangle+|100\rangle+|111\rangle}{2}$  and involves a similar procedure [49]. In all these instances, the measurement of the encoded Clifford operator is realized by applying transversally the operator conditioned on the qubits in a cat state.

We now show a general method that can be used to implement geometrically any conditional transversal Clifford operation with conditioning on the cat state. Let  $O$  be a Clifford gate acting on the first qubits from some set of blocks. As we discussed in Sec. III A 3, under this unitary the stabilizer and the gauge group transform in such a way that we can always find an element with an arbitrary combination of Pauli matrices on the first qubits. If we write this element in the form

$$\hat{G} = G_1 \otimes G_{2,\dots,n}, \quad (45)$$

where  $G_1$  is a tensor product of Pauli matrices acting on the first qubits from the blocks and  $G_{2,\dots,n}$  is an operator on the rest of the qubits, then applying  $O$  conditioned on the first qubit in a cat state transforms this stabilizer or gauge-group element as follows:

$$\begin{aligned} I^c \otimes G_1 \otimes G_{2,\dots,n} &= |0\rangle\langle 0|^c \otimes G_1 \otimes G_{2,\dots,n} \\ &\quad + |1\rangle\langle 1|^c \otimes G_1 \otimes G_{2,\dots,n} \\ &\rightarrow |0\rangle\langle 0|^c \otimes G_1 \otimes G_{2,\dots,n} \end{aligned}$$



$$+ |1\rangle\langle 1|^c \otimes OG_1O^\dagger \otimes G_{2,\dots,n}, \quad (46)$$

where the superscript  $c$  denotes the control qubit from the cat state. We can implement this operation by choosing the factor  $G_1$  to be the same as the one we would use if we wanted to implement the operation  $O$  according to the previously described procedure. Then we can apply the following Hamiltonian:

$$\hat{H}_{C(O)}(t) = -|0\rangle\langle 0|^c \otimes G_1 \otimes G_{2,\dots,n} - \alpha(t)|1\rangle\langle 1|^c \otimes H_O(t) \otimes G_{2,\dots,n}, \quad (47)$$

where  $-H_O(t) \otimes G_{2,\dots,n}$  is the Hamiltonian that we would use for the implementation of the operation  $O$  and  $\alpha(t)$  is a real parameter chosen such that at every moment the operator  $\alpha(t)|1\rangle\langle 1|^c \otimes H_O(t) \otimes G_{2,\dots,n}$  has the same instantaneous spectrum as the operator  $|0\rangle\langle 0|^c \otimes G_1 \otimes G_{2,\dots,n}$ . This guarantees that the overall Hamiltonian is degenerate and the geometric transformation of each of its eigenspaces is given by the operator

$$\hat{U}_g(t) = |0\rangle\langle 0|^c \otimes I_1 \otimes I_{2,\dots,n} + |1\rangle\langle 1|^c \otimes U_O(t) \otimes I_{2,\dots,n}, \quad (48)$$

where  $U_O(t)$  is the geometric transformation on the first qubits generated by  $-H_O(t) \otimes G_{2,\dots,n}$ . Since we presented the constructions of our basic Clifford operations up to an overall phase, the operation  $U_O(t)$  may differ from the desired operation by a phase. This phase can be corrected by applying a suitable gate on the control qubit from the cat state (we explain how this can be done in Sec. III C). We remark that a Hamiltonian of the type given in Eq. (47) requires fine tuning of the parameter  $\alpha(t)$  and generally can be complicated. Our goal in this section is to prove that universal fault-tolerant holonomic computation is possible in principle. In Sec. V we show that depending on the code one can find more natural implementations of these operations.

If we want to apply a second conditional Clifford operation  $Q$  on the first qubits in the blocks, we can do this as follows. Imagine that if we had to apply the operation  $Q$  following the operation  $O$ , we would use the Hamiltonian  $\hat{H}_Q(t) = -H_Q(t) \otimes G'_{2,\dots,n}$ , where  $\hat{H}_Q(0) = OG'_1O^\dagger \otimes G'_{2,\dots,n}$  is a suitable element of the stabilizer or the gauge group after the application of  $O$ . Before the application of  $O$ , that element would have had the form  $G'_1 \otimes G'_{2,\dots,n}$ . Under the application of a conditional  $O$ , the element  $G'_1 \otimes G'_{2,\dots,n}$  transforms to  $|0\rangle\langle 0|^c \otimes G'_1 \otimes G'_{2,\dots,n} + |1\rangle\langle 1|^c \otimes OG'_1O^\dagger \otimes G'_{2,\dots,n}$  which can be used (with a minus sign) as a starting Hamiltonian for a subsequent operation. In particular, we can implement the conditional  $Q$  following the conditional  $O$  using the Hamiltonian

$$\hat{H}_{C(Q)}(t) = -|0\rangle\langle 0|^c \otimes G'_1 \otimes G'_{2,\dots,n} - \beta(t)|1\rangle\langle 1|^c \otimes H_Q(t) \otimes G'_{2,\dots,n}, \quad (49)$$

where the factor  $\beta(t)$  guarantees that there is no splitting of the energy levels. Subsequent operations can be applied analogously. Using this general method, we can implement a

unitary whose geometric part is equal to any transversal Clifford operation conditioned on the cat state.

### C. Preparing and using the cat state

In addition to transversal operations, a complete fault-tolerant scheme requires the ability to prepare, verify, and use a special ancillary state such as the cat state  $(|00\dots 0\rangle + |11\dots 1\rangle)/\sqrt{2}$  proposed by Shor [1]. This can also be done using our geometric approach. Since the cat state is known and its construction is non-fault-tolerant, we can prepare it by simply treating each initially prepared qubit as a simple code [with  $\tilde{G}$  in Eq. (6) being trivial] and updating the stabilizer of the code via the applied geometric transformation as the operation progresses. The stabilizer of the prepared cat state is generated by  $Z_i Z_j$ ,  $i < j$ . Transversal unitary operations between the cat state and other codewords are applied as described in Secs. III A and III B.

We also have to be able to measure the parity of the state, which requires the ability to apply successive C-NOT operations from two different qubits in the cat state to the same ancillary qubit initially prepared in the state  $|0\rangle$ . We can regard a qubit in state  $|0\rangle$  as a simple code with stabilizer  $\langle Z \rangle$ , and we can apply the first C-NOT as described before. Even though after this operation the state of the target qubit is unknown, the second C-NOT gate can be applied via the same interaction since the transformation undergone by each eigenspace would still be equivalent to the desired C-NOT and at the end when we measure the qubit we project onto one of the eigenspaces.

### D. Fault tolerance of the scheme

We showed how we can generate any transversal operation on the code space geometrically, assuming that the state is nonerroneous. But what if an error occurs on one of the qubits?

At any moment, we can distinguish two types of errors—those that result in transitions between the ground and the excited spaces of the current Hamiltonian and those that result in transformations inside the eigenspaces. Due to the discretization of errors in QEC, it suffices to prove correctness for each type separately. The key property of our construction is that the transformation undergone by each of the eigenspaces is equivalent to the same transversal operation. Because of this, if we are applying a unitary on the first qubit, an error on that qubit will remain localized regardless of whether it causes an excitation or not. If the error occurs on one of the other qubits, at the end of the transformation the result would be the desired single-qubit unitary gate plus the error on the other qubit, which is correctable.

It is remarkable that even though the Hamiltonian couples qubits within the same block, single-qubit errors do not propagate. This is because the coupling between the qubits amounts to a change in the relative phase between the ground and excited spaces, but the latter is irrelevant since either it is equivalent to a gauge transformation or when we apply a correcting operation we project onto one of the eigenspaces. In the case of the C-NOT gate, an error can

propagate between the control and the target qubits, but it never results in two errors within the same codeword.

#### IV. EFFECTS ON THE ACCURACY THRESHOLD FOR ENVIRONMENTAL NOISE

Since the method we presented conforms completely to a given fault-tolerant scheme, it would not affect the error threshold per operation for that scheme. Some of its features, however, would affect the threshold for *environmental* noise.

First, observe that when applying the Hamiltonian in Eq. (8), we cannot at the same time apply operations on the other qubits on which the factor  $\tilde{G}$  acts nontrivially. Thus, some operations at the lowest level of concatenation that would otherwise be implemented simultaneously might have to be implemented serially. The effect of this is equivalent to slowing down the circuit by a constant factor. [Note that we could also vary the factor  $\tilde{G}$  simultaneously with  $H(t)$ , but in order to obtain the same precision as that we would achieve by a serial implementation, we would have to slow down the change in the Hamiltonian by the same factor.] The slowdown factor resulting from this loss of parallelism is usually small since this problem occurs only at the lowest level of concatenation. For example, when implementing encoded single-qubit gates with the Bacon-Shor code, we can apply operations on up to six out of the nine qubits in a block simultaneously. As we show in Sec. V, we can address any two qubits in a row or column using our method by taking  $\tilde{G}$  in Eq. (8) to be a single-qubit operator  $Z$  or  $X$  on the third qubit in the same row or column. The Hamiltonians used to apply operations on the two qubits commute with each other at all times and do not interfere. A similar phenomenon holds for the implementation of the encoded C-NOT gate or the operations involving the cat state. Thus, for the Bacon-Shor code we have a slowdown due to parallelism by a factor of 1.5.

A more significant slowdown results from the fact that the evolution is adiabatic. In order to obtain a rough estimate of the slowdown due specifically to the adiabatic requirement, we will compare the time  $T_h$  needed for the adiabatic implementation of a given gate with precision  $1 - \delta$  to the time  $T_d$  needed for a dynamical realization of the same gate with the same strength of the Hamiltonian. We will consider a realization of the  $X$  gate via the unitary interpolation [50]

$$\hat{H}(t) = -V_X[\tau(t)]ZV_X^\dagger[\tau(t)] \otimes \tilde{G},$$

$$V_X[\tau(t)] = \exp\left(i\tau(t)\frac{\pi}{2T_h}X\right), \quad (50)$$

where  $\tau(0)=0$ ,  $\tau(T_h)=T_h$ . Thus the energy gap of the Hamiltonian is constant. The optimal dynamical implementation of the same gate is via the Hamiltonian  $-X$  for time  $T_d=\frac{\pi}{2}$ .

As we argued in Sec. III, the accuracy with which the adiabatic approximation holds for the Hamiltonian in Eq. (50) is the same as that for the Hamiltonian

$$H(t) = V_X(\tau(t))ZV_X^\dagger(\tau(t)). \quad (51)$$

We now present estimates for two different choices of the function  $\tau(t)$ . The first one is

$$\tau(t) = t. \quad (52)$$

In this case the Schrödinger equation can be easily solved in the instantaneous eigenbasis of the Hamiltonian in Eq. (51). For the probability that the initial ground state remains a ground state at the end of the evolution, we obtain

$$p = \frac{1}{1 + \varepsilon^2} + \frac{\varepsilon^2}{1 + \varepsilon^2} \cos^2\left(\frac{\pi}{4\varepsilon}\sqrt{1 + \varepsilon^2}\right) = 1 - \delta, \quad (53)$$

where

$$\varepsilon = \frac{T_d}{T_h}. \quad (54)$$

Expanding in powers of  $\varepsilon$  and averaging the square of the cosine whose period is much smaller than  $T_h$ , we obtain the condition

$$\varepsilon^2 \leq 2\delta. \quad (55)$$

Assuming, for example, that  $\delta \approx 10^{-4}$  (approximately the threshold for the nine-qubit Bacon-Shor code [40]), we obtain that the time of evolution for the adiabatic case must be about 70 times longer than that in the dynamical case.

It is known, however, that if  $H(t)$  is smooth and all its derivatives vanish at  $t=0$  and  $t=T_h$ , the adiabatic error decreases superpolynomially with  $T_h$  [47]. To achieve this, we will choose

$$\tau(t) = \frac{1}{a} \int_0^t dt' e^{-1/\sin(\pi t'/T_h)}, \quad a = \int_0^{T_h} dt' e^{-1/\sin(\pi t'/T_h)}. \quad (56)$$

For this interpolation, by a numerical solution we obtain that when  $T_h/T_d \approx 17$  the error is already of the order of  $10^{-6}$ , which is well below the threshold values obtained for the Bacon-Shor codes [40]. This is a remarkable improvement in comparison to the previous interpolation which shows that the smoothness of the Hamiltonian plays an important role in the performance of the scheme.

An additional slowdown in comparison to a perfect dynamical scheme may result from the fact that the constructions for some of the standard gates we presented involve long sequences of loops. With more efficient parameter paths, however, it should be possible to reduce this slowdown to minimum. An approach for finding optimal loops presented in Ref. [51] may be useful in this respect.

In comparison to a dynamical implementation, the allowed rate of environmental noise for the holonomic case would decrease by a factor similar to the slowdown factor. In practice, however, dynamical gates are not perfect and the holonomic approach may be advantageous if it gives rise to a better operational precision.

We finally point out that an error in the factor  $H(t)$  in the Hamiltonian in Eq. (8) would result in an error on the first qubit according to Eq. (21). Such an error clearly has to be below the accuracy threshold. More dangerous errors, how-

ever, are also possible. For example, if the degeneracy of the Hamiltonian is broken, this can result in an unwanted dynamical transformation affecting all qubits on which the Hamiltonian acts nontrivially. Such multiqubit errors have to be of higher order in the threshold, which imposes more severe restrictions on the Hamiltonian.

### V. FAULT-TOLERANT HOLONOMIC COMPUTATION WITH LOW-WEIGHT HAMILTONIANS

The weight of the Hamiltonians needed for the scheme we described depends on the weight of the stabilizer or gauge-group elements. Remarkably, certain codes possess stabilizer or gauge-group elements of low weight covering all qubits in the code, which allows us to perform holonomic computation using low-weight Hamiltonians. Here we will consider as an example a subsystem generalization of the nine-qubit Shor code [9]—the Bacon-Shor code [26,27]—which has particularly favorable properties for fault-tolerant computation [40,52]. In the nine-qubit Bacon-Shor code, the gauge group is generated by the weight-two operators  $Z_{k,j}Z_{k,j+1}$  and  $X_{j,k}X_{j+1,k}$ , where the subscripts label the qubits by row and column when they are arranged in a  $3 \times 3$  square lattice. Since the Bacon-Shor code is a CSS code, the C-NOT gate has a direct transversal implementation. We now show that the C-NOT gate can be realized using at most weight-three Hamiltonians.

If we want to apply a C-NOT gate between two qubits each of which is, say, in the first row and column of its block, we can use as a starting Hamiltonian  $-Z_{1,1}^t \otimes Z_{1,2}^t$ , where the superscript  $t$  signifies that these are operators in the target block. We can then apply the C-NOT gate as described in Sec. III. After the operation, however, this gauge-group element will transform to  $-Z_{1,1}^t \otimes Z_{1,1}^c \otimes Z_{1,2}^t$ . If we now want to implement a C-NOT gate between the qubits with index  $\{1,2\}$  using as a starting Hamiltonian the operator  $-Z_{1,1}^t \otimes Z_{1,1}^c \otimes Z_{1,2}^t$  according to the same procedure, we will have to use a four-qubit Hamiltonian. Of course, at this point we can use the starting Hamiltonian  $-Z_{1,2}^t \otimes Z_{1,3}^t$ , but if we had also applied a C-NOT between the qubits labeled  $\{1,3\}$ , this operator would not be available—it would have transformed to  $-Z_{1,2}^t \otimes Z_{1,3}^c \otimes Z_{1,3}^t$ .

What we can do instead is to use as a starting Hamiltonian the operator  $-Z_{1,1}^t \otimes Z_{1,2}^t \otimes Z_{1,2}^c$  which is obtained from the gauge-group element  $Z_{1,1}^t \otimes Z_{1,1}^c \otimes Z_{1,2}^t \otimes Z_{1,2}^c$  after the application of the C-NOT between the qubits with index  $\{1,1\}$ . Since the C-NOT gate is its own inverse, we can regard the factor  $Z_{1,1}^t$  as  $\tilde{G}$  in Eq. (44) and use this starting Hamiltonian to apply our procedure backward. Thus we can implement any transversal C-NOT gate using at most weight-three Hamiltonians.

Since the encoded  $X$ ,  $Y$ , and  $Z$  operations have a bitwise implementation, we can always apply them according to our procedure using Hamiltonians of weight 2. For the Bacon-Shor code, the encoded Hadamard gate can be applied via bitwise Hadamard transformations followed by a rotation of the grid by a  $90^\circ$  angle [40]. The encoded phase gate can be implemented by using the encoded C-NOT and an ancilla.

The preparation and measurement of the cat state can also be done using Hamiltonians of weight 2. To prepare the cat state, we first prepare all qubits in the state  $(|0\rangle + |1\rangle)/\sqrt{2}$ , which can be done by measuring each of them in the  $\{|0\rangle, |1\rangle\}$  basis (this ability is assumed for any type of computation) and applying the transformation  $-Z \rightarrow -X$  or  $Z \rightarrow -X$  depending on the outcome. To complete the preparation of the cat state, apply a two-qubit transformation between the first qubit and each of the other qubits ( $j > 1$ ) via the transformation

$$-I_1 \otimes X_j \rightarrow -Z_1 \otimes Z_j. \quad (57)$$

Single-qubit transformations on qubits from the cat state can be applied according to the method described in Sec. III using at most weight-two Hamiltonians.

To measure the parity of the state, we need to apply successively C-NOT operations from two different qubits in the cat state to the same ancillary qubit initially prepared in the state  $|0\rangle$ . As described in Sec. III, this can also be done according to our method and requires Hamiltonians of weight 2.

For universal computation with the Bacon-Shor code, we also need to be able to apply one encoded transformation outside of the Clifford group. As we mentioned earlier, in order to implement the Toffoli gate or the  $\pi/8$  gate, it is sufficient to be able to implement a C-NOT gate conditioned on a cat state. For the Bacon-Shor code, the C-NOT gate has a transversal implementation, so the conditioned C-NOT gate can be realized by a series of transversal Toffoli operations between the cat state and the two encoded states. We now show that this gate can be implemented using at most three-qubit Hamiltonians.

Nielsen and Chuang [33] provide a circuit for implementing the Toffoli gate as a sequence of one- and two-qubit gates. We will use the same circuit, except that we flip the control and target qubits in every C-NOT gate using the identity

$$(R_1 \otimes R_2)C_{1,2}(R_1 \otimes R_2) = C_{2,1}, \quad (58)$$

where  $R_i$  denotes a Hadamard gate on the qubit labeled by  $i$  and  $C_{i,j}$  denotes a C-NOT gate between qubits  $i$  and  $j$  with  $i$  being the control and  $j$  being the target. Let  $\text{Toffoli}_{i,j,k}$  denote the Toffoli gate on qubits  $i$ ,  $j$ , and  $k$ , with  $i$  and  $j$  being the two control qubits and  $k$  being the target qubit, and let  $S_i$  and  $T_i$  denote the phase and  $\pi/8$  gates on qubit  $i$ , respectively. Then the Toffoli gate on three qubits (the first one of which we will assume to belong to the cat state) can be written as

$$\begin{aligned} \text{Toffoli}_{1,2,3} = & R_2 C_{3,2} R_3 T_3^\dagger R_3 R_1 C_{3,1} R_3 T_3 R_3 C_{3,2} R_3 T_3^\dagger R_3 C_{3,1} \\ & \times R_3 T_3 R_3 R_2 T_2^\dagger R_2 C_{2,1} R_2 T_2^\dagger R_2 C_{2,1} R_2 S_2 R_1 T_1. \end{aligned} \quad (59)$$

To show that each of the above gates can be implemented according to our geometric approach using Hamiltonians of weight at most 3, we will need an implementation of the C-NOT gate which is suitable for the case when we have a stabilizer or gauge-group element of the form

$$\hat{G} = X \otimes \tilde{G}, \quad (60)$$

where the factor  $X$  acts on the target qubit and  $\tilde{G}$  acts trivially on the control qubit. By a similar argument to the one in Sec. III, one can verify that in this case the C-NOT gate can be implemented as follows: apply the operation  $S^\dagger$  on the control qubit (we describe how to do this for our particular case below) together with the transformation

$$-X \otimes \tilde{G} \rightarrow -Z \otimes \tilde{G} \rightarrow X \otimes \tilde{G} \quad (61)$$

on the target qubit, followed by the transformation

$$\begin{aligned} I^c \otimes X \otimes \tilde{G} &\rightarrow -(|0\rangle\langle 0|^c \otimes Z + |1\rangle\langle 1|^c \otimes Y) \otimes \tilde{G} \\ &\rightarrow -I^c \otimes X \otimes \tilde{G}. \end{aligned} \quad (62)$$

Since the second and the third qubits belong to blocks encoded with the Bacon-Shor code, there are weight-two elements of the initial gauge group of the form  $Z \otimes Z$  covering all qubits. The stabilizer generators on the cat state are also of this type. Following the transformation of these operators according to the sequence of operations in Eq. (59), one can see that before every C-NOT gate in this sequence, there is an element of the form given in Eq. (60) with  $\tilde{G}=Z$  that can be used to implement the C-NOT gate as described, provided that we can implement the gate  $S^\dagger$  on the control qubit. We also point out that all single-qubit operations on qubit 1 in this sequence can be implemented according to the procedure described in Sec. III, since at every step we have a weight-two stabilizer element on that qubit with a suitable form. Therefore, all we need to show is how to implement the necessary single-qubit operations on qubits 2 and 3. Due to the complicated transformation of the gauge-group elements during the sequence of operations [Eq. (59)], we introduce a method of applying a single-qubit operation with a starting Hamiltonian that acts trivially on the qubit. For implementing single-qubit operations on qubits 2 and 3 we use as a starting Hamiltonian the operator

$$\hat{H}(0) = -I_i \otimes X_1 \otimes \tilde{Z}, \quad i=2,3, \quad (63)$$

where the first factor ( $I_i$ ) acts on the qubit on which we want to apply the operation (2 or 3) and  $X_1 \otimes \tilde{Z}$  is the transformed (after the Hadamard gate  $R_1$ ) stabilizer element of the cat state that acts nontrivially on qubit 1 (the factor  $\tilde{Z}$  acts on some other qubit in the cat state).

To implement a single-qubit gate on qubit 3, for example, we first apply the interpolation

$$-I_3 \otimes X_1 \otimes \tilde{Z} \rightarrow -Z_3 \otimes X_1 \otimes \tilde{Z}. \quad (64)$$

This results in a two-qubit geometric transformation  $U_{1,3}$  on qubits 1 and 3. We do not have to calculate this transformation exactly since we will undo it later, but the fact that each eigenspace undergoes the same two-qubit geometric transformation can be verified similarly to the C-NOT gate we described in Sec. III.

At this point, the Hamiltonian is of the form (7) with respect to qubit 3, and we can apply any single-qubit unitary gate  $V_3$  according to the method described in Sec. III. This

transforms the Hamiltonian to  $-V_3 Z_3 V_3^\dagger \otimes X_1 \otimes \tilde{Z}$ . We can now “undo” the transformation  $U_{1,3}$  by the interpolation

$$-V_3 Z_3 V_3^\dagger \otimes X_1 \otimes \tilde{Z} \rightarrow -I_3 \otimes X_1 \otimes \tilde{Z}. \quad (65)$$

The latter transformation is the inverse of Eq. (64) up to the single-qubit unitary transformation  $V_3$ , i.e., it results in the transformation  $V_3 U_{1,3}^\dagger V_3^\dagger$ . Thus the net result is

$$V_3 U_{1,3}^\dagger V_3^\dagger V_3 U_{1,3} = V_3, \quad (66)$$

which is the desired single-qubit unitary transformation on qubit 3. We point out that during this transformation, a single-qubit error can propagate between qubits 1 and 3, but this is not a problem since we are implementing a transversal Toffoli operation and such an error would not result in more than one error per block of the code.

We showed that for the Bacon-Shor code our scheme can be implemented with at most three-local Hamiltonians. This is optimal for the construction we presented since there are no nontrivial codes with stabilizer or gauge-group elements of weight smaller than 2 covering all qubits. One could argue that since the only Hamiltonians that leave the code space invariant are superpositions of elements of the stabilizer or the gauge group, one cannot do better than this. However, it may be possible to approximate the necessary Hamiltonians with sufficient precision using two-local interactions. A possible direction to consider in this respect is the gadget techniques introduced in Ref. [53] and developed further in Refs. [54,55]. This is left as a problem for future investigation.

## VI. CONCLUSION

We described a scheme for fault-tolerant holonomic computation on stabilizer codes, which demonstrates that HQC is a scalable method of computation. The scheme opens the possibility for combining the software protection of error correction with the inherent robustness of HQC against control imperfections. Our construction uses Hamiltonians that are elements of the stabilizer or the gauge group of the code and works by adiabatically varying these Hamiltonians in a manner which generates unitaries whose geometric parts are equal to transversal operations from the point of view of the basis of the full Hilbert space. By composing these transversal operations as in a given standard dynamical fault-tolerant scheme, we thus transport the code space (or the redundant subspaces that constitute a code subsystem) adiabatically around loops that give rise to holonomies equal to encoded operations inside the code space. The Hamiltonians needed for implementing two-qubit gates are at least three-local. We showed that computation with at most three-local Hamiltonians is possible with the Bacon-Shor code.

It is interesting to point out that the adiabatic regime in which our scheme operates is consistent with the model of Markovian decoherence. In Ref. [56] it was argued that the standard dynamical paradigm of fault tolerance is based on assumptions that are in conflict with the rigorous derivation of the Markovian limit. Although the threshold theorem has been extended to non-Markovian models [57–59], the Markovian assumption is an accurate approximation for a wide

range of physical scenarios [60]. It also allows for a much simpler description of the evolution in comparison to non-Markovian models (see, e.g., Ref. [61]). In Ref. [56] it was shown that the weak-coupling-limit derivation of the Markovian approximation is consistent with computational methods that employ slow transformations, such as adiabatic quantum computation [62] or HQC. A theory of fault tolerance for the adiabatic model of computation at present is not known although some steps in this direction have been undertaken [63,64]. Our hybrid HQC-QEC scheme provides a solution for the case of HQC. We point out, however, that it is an open problem whether the Markovian approximation makes sense for a fixed value of the adiabatic slowness parameter when the circuit increases in size. Giving a definitive answer to this question requires a rigorous analysis of the accumulation of non-Markovian errors due to deviation from perfect adiabaticity.

Applying the present strategy to actual physical systems might require modifying our abstract construction in accordance with the available interactions, possibly using linear combinations of stabilizer or gauge-group elements rather than single elements as the basic Hamiltonians. Given that simple QEC codes and two-qubit geometric transformations have been realized using NMR [65,66] and ion-trap [67,68] techniques, these systems seem particularly suitable for hybrid HQC-QEC implementations.

We hope that the techniques presented in this paper might prove useful in other areas as well. It is possible that some combination of transversal adiabatic transformations and active correction could provide a solution to the problem of fault tolerance in the adiabatic model of computation as well.

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#### APPENDIX A: PROOF OF PROPOSITION 2

To prove proposition 2, observe that the Hermitian unitary matrix

$$W^\theta = \begin{pmatrix} 0 & ie^{-i\theta} \\ -ie^{i\theta} & 0 \end{pmatrix} \quad (\text{A1})$$

has the properties

$$[W^\theta, V^{\theta\pm}] = 0, \quad (\text{A2})$$

$$\{W^\theta, Z\} = 0, \quad (\text{A3})$$

where  $[\cdot, \cdot]$  and  $\{\cdot, \cdot\}$  denote a commutator and an anticommutator, respectively. This means that

$$\begin{aligned} W^\theta H^{\theta\pm}(t) W^\theta &= W^\theta [f(t)Z + g(t)V^{\theta\pm}ZV^{\theta\mp}] W^\theta \\ &= -f(t)Z - g(t)V^{\theta\pm}ZV^{\theta\mp} \\ &= -H^{\theta\pm}(t), \end{aligned} \quad (\text{A4})$$

i.e., the unitary  $W^\theta$  flips the ground and excited spaces of  $H^{\theta\pm}(t)$  for any  $t$ .

The unitaries  $U_{0,1}^{\theta\pm}$ , given by Eq. (13) for  $H(t)=H^{\theta\pm}(t)$ , are therefore related by

$$U_0^{\theta\pm}(t) = W^\theta U_1^{\theta\pm}(t) W^\theta. \quad (\text{A5})$$

Using the fact that

$$\begin{aligned} W^\theta|0\rangle &= -ie^{i\theta}|1\rangle, \\ W^\theta|1\rangle &= ie^{-i\theta}|0\rangle, \end{aligned} \quad (\text{A6})$$

from Eqs. (14) and (15) one can see that Eq. (A5) implies

$$U_{A_0}^{\theta\pm}(t) = W^\theta U_{A_1}^{\theta\pm}(t) W^\theta. \quad (\text{A7})$$

Let us define the eigenstates of  $H^{\theta\pm}(t)$  at time  $T$  as  $|\phi_0^{\theta\pm}(T)\rangle = V^{\theta\pm}|0\rangle$  and  $|\phi_1^{\theta\pm}(T)\rangle = V^{\theta\pm}|1\rangle$ . Expression (15) can then be written as

$$\begin{aligned} U_{A_0}^{\theta\pm}(T) &= e^{i\alpha_0^{\theta\pm}} V^{\theta\pm}|0\rangle\langle 0|, \\ U_{A_1}^{\theta\pm}(T) &= e^{i\alpha_1^{\theta\pm}} V^{\theta\pm}|1\rangle\langle 1|, \end{aligned} \quad (\text{A8})$$

where  $\alpha_0^{\theta\pm}$  and  $\alpha_1^{\theta\pm}$  are geometric phases. Without explicitly calculating the geometric phases, from Eqs. (A1) and (A6)–(A8), we obtain

$$\begin{aligned} e^{i\alpha_0^{\theta\pm}} V^{\theta\pm}|0\rangle\langle 0| &= e^{i\alpha_1^{\theta\pm}} W^\theta V^{\theta\pm}|1\rangle\langle 1| W^\theta \\ &= e^{i\alpha_1^{\theta\pm}} V^{\theta\pm} W^\theta|1\rangle\langle 1| W^\theta \\ &= e^{i\alpha_1^{\theta\pm}} V^{\theta\pm}|0\rangle\langle 0|, \end{aligned} \quad (\text{A9})$$

i.e.,

$$e^{i\alpha_0^{\theta\pm}} = e^{i\alpha_1^{\theta\pm}}. \quad (\text{A10})$$

Therefore, up to a global phase, Eq. (21) yields

$$U^{\theta\pm}(T) \sim V^{\theta\pm}. \quad (\text{A11})$$

#### APPENDIX B: CALCULATING THE HOLONOMY FOR THE Z GATE

##### 1. Linear interpolation

We first demonstrate how to calculate the ground-space holonomy for the Z gate for the case of linear interpolations along the segments of the path, i.e., when  $f(t)$  and  $g(t)$  in Eq. (24) are

$$f(t) = 1 - \frac{t}{T}, \quad g(t) = \frac{t}{T}. \quad (\text{B1})$$

In order to calculate the holonomy in Eq. (4) corresponding to our construction of the Z gate, we need to define a

single-valued orthonormal basis of the ground space of the Hamiltonian along the loop described by Eq. (27). Since the Hamiltonian has the form given in Eq. (10) at all times, it is convenient to choose the basis of the form

$$|jk; \lambda\rangle = |\chi_j(\lambda)\rangle |\tilde{\psi}_{jk}\rangle, \quad (B2)$$

$$j = 0, 1; \quad k = 1, \dots, 2^{n-2},$$

where  $|\chi_0(\lambda(t))\rangle$  and  $|\chi_1(\lambda(t))\rangle$  are ground and excited states of  $H(t)$  and  $|\tilde{\psi}_{0k}\rangle$  and  $|\tilde{\psi}_{1k}\rangle$  are fixed orthonormal bases of the subspaces that support the projectors  $\tilde{P}_0$  and  $\tilde{P}_1$  defined in Eq. (11), respectively. The eigenstates  $|\chi_0(\lambda(t))\rangle$  and  $|\chi_1(\lambda(t))\rangle$  are defined up to an overall phase, but we have to choose the phase such that the states are single-valued along the loop.

Observe that because of this choice of basis, the matrix elements given in Eq. (3) become

$$(A_\mu)_{jk, j'k'} = \langle jk; \lambda | \frac{\partial}{\partial \lambda^\mu} | j'k'; \lambda \rangle$$

$$= \langle \chi_j(\lambda) | \frac{\partial}{\partial \lambda^\mu} | \chi_{j'}(\lambda) \rangle \langle \tilde{\psi}_{jk} | \tilde{\psi}_{j'k'} \rangle$$

$$= \langle \chi_j(\lambda) | \frac{\partial}{\partial \lambda^\mu} | \chi_{j'}(\lambda) \rangle \delta_{jj'} \delta_{kk'}, \quad (B3)$$

i.e., the matrix  $A_\mu$  is diagonal. (Since we are looking only at the ground space, we are not writing the index of the energy level.) We can therefore drop the path-ordering operator. The resulting unitary matrix  $U_{jk, j'k'}^\gamma$  acting on the subspace spanned by  $\{|jk; \lambda(0)\rangle\}$  is also diagonal and its diagonal elements are

$$U_{jk, jk}^\gamma = \exp\left(\oint_\gamma \langle \chi_j(\lambda) | \frac{\partial}{\partial \lambda^\mu} | \chi_j(\lambda) \rangle d\lambda^\mu\right). \quad (B4)$$

These are precisely the Berry phases for the loops described by the states  $|\chi_j(\lambda)\rangle$ . Since the loop in parameter space consists of four line segments, we can write the last expression as

$$U_{jk, jk}^\gamma = \exp\left(\sum_{i=1}^4 \int_{\gamma_i} \langle \chi_j(\lambda) | \frac{\partial}{\partial \lambda^\mu} | \chi_j(\lambda) \rangle d\lambda^\mu\right), \quad (B5)$$

where  $\gamma_i$ ,  $i=1, 2, 3, 4$ , are the segments indexed by their order corresponding to Eq. (27). If we parametrize each line segment by the dimensionless time  $0 \leq s \leq 1$ , we obtain

$$U_{jk, jk}^\gamma = \exp\left(\sum_{i=1}^4 \int_0^1 \langle \chi_j^i(s) | \frac{d}{ds} | \chi_j^i(s) \rangle ds\right), \quad (B6)$$

where the superscript  $i$  in  $|\chi_j^i(s)\rangle$  indicates the segment. In the  $\{|0\rangle, |1\rangle\}$  basis, we will write these states as

$$|\chi_j^i(s)\rangle = \begin{pmatrix} a_j^i(s) \\ b_j^i(s) \end{pmatrix}, \quad j = 0, 1, \quad i = 1, 2, 3, 4, \quad (B7)$$

where  $|a_j^i(s)|^2 + |b_j^i(s)|^2 = 1$ .

Along the segment  $\gamma_1$ , the Hamiltonian has the form  $\hat{H}(s) = H_1(s) \otimes \tilde{P}_0 - H_1(s) \otimes \tilde{P}_1$ , where

$$H_1(t) = (1-s)Z + sX, \quad (B8)$$

i.e., the states  $|\chi_0^1(s)\rangle$  and  $|\chi_1^1(s)\rangle$  are the ground and excited states of  $H_1(s)$ . For these states we obtain

$$a_0^1(s) = \frac{(1-s + \sqrt{1-2s+2s^2})e^{i\omega_0^1(s)}}{\sqrt{2-4s+4s^2 + (2-2s)\sqrt{1-2s+2s^2}}}, \quad (B9)$$

$$b_0^1(s) = \frac{se^{i\omega_0^1(s)}}{\sqrt{2-4s+4s^2 + (2-2s)\sqrt{1-2s+2s^2}}}, \quad (B10)$$

$$a_1^1(s) = \frac{(1-s - \sqrt{1-2s+2s^2})e^{i\omega_1^1(s)}}{\sqrt{2-4s+4s^2 - (2-2s)\sqrt{1-2s+2s^2}}}, \quad (B11)$$

$$b_1^1(s) = \frac{se^{i\omega_1^1(s)}}{\sqrt{2-4s+4s^2 - (2-2s)\sqrt{1-2s+2s^2}}}, \quad (B12)$$

where  $\omega_j^1(s)$  are arbitrary phases which have to be chosen so that when we complete the loop, the phases of the corresponding states will return to their initial values mod  $2\pi$ . We will define the loops as interpolating between the following intermediate states defined with their overall phases:

$$|\psi_0(\lambda)\rangle: |0\rangle \rightarrow |f_+^0\rangle \rightarrow |1\rangle \rightarrow |f_-^{\pi/2}\rangle \rightarrow |0\rangle, \quad (B13)$$

$$|\psi_1(\lambda)\rangle: |1\rangle \rightarrow |f_-^0\rangle \rightarrow |0\rangle \rightarrow |f_+^{\pi/2}\rangle \rightarrow |1\rangle, \quad (B14)$$

where

$$|f_\pm^\theta\rangle = \frac{|0\rangle \pm e^{i\theta}|1\rangle}{\sqrt{2}}. \quad (B15)$$

In other words, we impose the conditions  $|\chi_{0,1}^1(0)\rangle = |0, 1\rangle$ ,  $|\chi_{0,1}^1(1)\rangle = |f_\pm^0\rangle = |\chi_{0,1}^2(0)\rangle$ ,  $|\chi_{0,1}^2(1)\rangle = |1, 0\rangle = |\chi_{0,1}^3(0)\rangle$ ,  $|\chi_{0,1}^3(1)\rangle = |f_\mp^{\pi/2}\rangle = |\chi_{0,1}^4(0)\rangle$ , and  $|\chi_{0,1}^4(1)\rangle = |0, 1\rangle$ .

From Eqs. (B9) and (B10) we see that  $a_0^1(0) = e^{i\omega_0^1(0)}$ ,  $b_0^1(0) = 0$  and  $a_0^1(1) = \frac{1}{\sqrt{2}}e^{i\omega_0^1(1)}$ ,  $b_0^1(1) = \frac{1}{\sqrt{2}}e^{i\omega_0^1(1)}$ , so we can choose

$$\omega_0^1(s) = 0, \quad \forall s \in [0, 1]. \quad (B16)$$

Similarly, from Eqs. (B11) and (B12) it can be seen that  $a_1^1(0) = 0$ ,  $b_1^1(0) = e^{i\omega_1^1(0)}$  and  $a_1^1(1) = -\frac{1}{\sqrt{2}}e^{i\omega_1^1(1)}$ ,  $b_1^1(1) = \frac{1}{\sqrt{2}}e^{i\omega_1^1(1)}$ . This means that  $\omega_1^1(s)$  has to satisfy  $e^{i\omega_1^1(0)} = 1$ ,  $e^{i\omega_1^1(1)} = -1$ . We can choose any differentiable  $\omega_1^1(s)$  that satisfies

$$\omega_1^1(0) = 0, \quad \omega_1^1(1) = \pi. \quad (B17)$$

In order to calculate  $\int_0^1 \langle \chi_j^1(s) | \frac{d}{ds} | \chi_j^1(s) \rangle ds$ , we also need

$$\frac{d}{ds} |\chi_j^1(s)\rangle = \begin{pmatrix} \frac{d}{ds} a_j^1(s) \\ \frac{d}{ds} b_j^1(s) \end{pmatrix}. \quad (B18)$$

Differentiating Eqs. (B9)–(B12) yields

$$\frac{d}{ds}a_0^1(s) = -\frac{s(1-s+\sqrt{1-2s+2s^2})}{2\sqrt{2-4s+4s^2}[1-2s+2s^2+(1-s)\sqrt{1-2s+2s^2}]^{3/2}}, \quad (\text{B19})$$

$$\frac{d}{ds}b_0^1(s) = \frac{2-4s+3s^2+(2-2s)\sqrt{1-2s+2s^2}}{2\sqrt{2-4s+4s^2}[1-2s+2s^2+(1-s)\sqrt{1-2s+2s^2}]^{3/2}}, \quad (\text{B20})$$

$$\frac{d}{ds}a_1^1(s) = -\frac{s(1-s-\sqrt{1-2s+2s^2})e^{i\omega_1^1(s)}}{2\sqrt{2-4s+4s^2}[1-2s+2s^2-(1-s)\sqrt{1-2s+2s^2}]^{3/2}} + a_1^1(s)i\frac{d}{ds}\omega_1^1(s), \quad (\text{B21})$$

$$\frac{d}{ds}b_0^1(s) = -\frac{[2-4s+3s^2-(2-2s)\sqrt{1-2s+2s^2}]e^{i\omega_1^1(s)}}{2\sqrt{2-4s+4s^2}[1-2s+2s^2-(1-s)\sqrt{1-2s+2s^2}]^{3/2}} + b_1^1(s)i\frac{d}{ds}\omega_1^1(s). \quad (\text{B22})$$

By a straightforward substitution, we obtain

$$\langle \chi_0^1(s) | \frac{d}{ds} | \chi_0^1(s) \rangle = a_0^{1*}(s) \frac{d}{ds} a_0^1(s) + b_0^{1*}(s) \frac{d}{ds} b_0^1(s) = 0, \quad (\text{B23})$$

$$\langle \chi_1^1(s) | \frac{d}{ds} | \chi_1^1(s) \rangle = a_1^{1*}(s) \frac{d}{ds} a_1^1(s) + b_1^{1*}(s) \frac{d}{ds} b_1^1(s) = i \frac{d}{ds} \omega_1^1(s). \quad (\text{B24})$$

Thus the integrals are

$$\int_0^1 \langle \chi_0^1(s) | \frac{d}{ds} | \chi_0^1(s) \rangle ds = 0, \quad (\text{B25})$$

$$\int_0^1 \langle \chi_1^1(s) | \frac{d}{ds} | \chi_1^1(s) \rangle ds = i \omega_1^1(s) \Big|_0^1 = i\pi. \quad (\text{B26})$$

In the same manner, we calculate the contributions of the other three line segments. The results are

$$\int_0^1 \langle \chi_0^2(s) | \frac{d}{ds} | \chi_0^2(s) \rangle ds = 0, \quad (\text{B27})$$

$$\int_0^1 \langle \chi_1^2(s) | \frac{d}{ds} | \chi_1^2(s) \rangle ds = 0, \quad (\text{B28})$$

$$\int_0^1 \langle \chi_0^3(s) | \frac{d}{ds} | \chi_0^3(s) \rangle ds = i\frac{\pi}{2}, \quad (\text{B29})$$

$$\int_0^1 \langle \chi_1^3(s) | \frac{d}{ds} | \chi_1^3(s) \rangle ds = 0, \quad (\text{B30})$$

$$\int_0^1 \langle \chi_0^4(s) | \frac{d}{ds} | \chi_0^4(s) \rangle ds = 0, \quad (\text{B31})$$

$$\int_0^1 \langle \chi_1^4(s) | \frac{d}{ds} | \chi_1^4(s) \rangle ds = i\frac{\pi}{2}. \quad (\text{B32})$$

Putting everything together, for the diagonal elements of the holonomy we obtain

$$U_{0k,0k}^\gamma = e^{i(\pi/2)},$$

$$U_{1k,1k}^\gamma = e^{i(3\pi/2)}. \quad (\text{B33})$$

The holonomy transforms any state in the ground space of the initial Hamiltonian as

$$U^\gamma \sum_{jk} \alpha_{jk} |j\rangle | \tilde{\psi}_{jk} \rangle = e^{i(\pi/2)} \sum_{jk} (-1)^j \alpha_{jk} |j\rangle | \tilde{\psi}_{jk} \rangle, \quad j=0,1. \quad (\text{B34})$$

From the point of view of the full Hilbert space, this is effectively a  $Z$  gate on the first qubit up to an overall phase.

Note that other single-qubit transformations such as the Hadamard or the  $X$  gates, under which the eigenspaces of the Hamiltonian do not follow complete loops, can be obtained in a similar fashion by calculating the open-path expression in Eq. (2). In principle, the result of that calculation depends on the choice of basis  $\{|\alpha; \lambda\rangle\}$  which is defined up to a unitary gauge transformation. However, if the final eigenspace has a nonzero overlap with the initial one, this ambiguity can be removed by defining the frame in the final eigenspace to be the one which is “most parallel” to the initial frame [29]. In the case of the single-qubit Hadamard gate, there is a nonzero overlap between the initial and final eigenspaces, and the set of initial basis states  $|0k; 0\rangle = |0\rangle | \tilde{\psi}_{0k} \rangle$ ,  $|1k; 0\rangle = |1\rangle | \tilde{\psi}_{1k} \rangle$ ,  $k=1, \dots, 2^{n-2}$ , can be seen to be most parallel to the set of final basis states  $|0k; 1\rangle = |+\rangle | \tilde{\psi}_{0k} \rangle$ ,  $|1k; 1\rangle = |-\rangle | \tilde{\psi}_{1k} \rangle$ ,  $k=1, \dots, 2^{n-2}$ , respectively (for a precise definition of most parallel see Ref. [29]). Thus, the resulting open-path holonomy corresponds to flipping the phase of half of the basis vectors. For the single-qubit  $X$  gate, however, the final ground (excited) space is orthogonal to the initial ground (excited) space. Thus a gauge-invariant expression for the geometric transformation cannot be defined in this way.

We emphasize that expression (2) for the open-path geometric transformation taking place inside each eigenspace of the Hamiltonian, whether gauge invariant or not, is not the same as the geometric transformation being realized inside the code space which is what eventually gives rise to the logical transformation. For example, the single-qubit Z gate can be understood as a closed-loop holonomy in the eigenspaces of the Hamiltonian  $-Z \otimes \tilde{G}$ , but under this transformation the code space does not follow a loop; in fact, it becomes orthogonal to the initial code space so that a gauge-invariant expression for the geometric transformation taking place inside it cannot be defined. However, if a sequence of single-qubit Z gates on all qubits yields, say, an encoded Z gate (as in the case of CSS codes), then after we complete such a sequence we will obtain a nontrivial closed-loop holonomy in the code space which is equal to the encoded Z.

### 2. Unitary interpolation

The calculation is simpler if we choose a unitary interpolation,

$$f(t) = \cos \frac{\pi t}{2T}, \quad g(t) = \sin \frac{\pi t}{2T}. \quad (\text{B35})$$

Such an interpolation corresponds to a rotation of the Bloch sphere around a particular axis for each of the segments of the loop. The first two segments of the loop given in Eq. (27) are realized via the Hamiltonian

$$\hat{H}_{1,2}(t) = -V_Y^\dagger(t) Z V_Y(t) \otimes \tilde{G}, \quad V_Y(t) = \exp\left(it \frac{\pi}{2T} Y\right), \quad (\text{B36})$$

applied for time  $T$ , and the third and fourth segments are realized via the Hamiltonian

$$\hat{H}_{3,4}(t) = -V_X(t) Z V_X^\dagger(t) \otimes \tilde{G}, \quad V_X(t) = \exp\left(it \frac{\pi}{2T} X\right), \quad (\text{B37})$$

again applied for time  $T$ . Let us define the eigenstates of the Hamiltonian along the first two segments as

$$|\chi_0^{1,2}(t)\rangle = V_X(t)|0\rangle, \quad |\chi_1^{1,2}(t)\rangle = V_X(t)|1\rangle, \quad 0 \leq t \leq T \quad (\text{B38})$$

and along the third and fourth segments as

$$|\chi_0^{3,4}(t)\rangle = -iV_Y^\dagger(t)Y|0\rangle, \\ |\chi_1^{3,4}(t)\rangle = -iV_Y^\dagger(t)Y|1\rangle, \quad 0 \leq t \leq T. \quad (\text{B39})$$

Notice that

$$|\chi_0^{1,2}(T)\rangle = -iY|0\rangle = |\chi_0^{3,4}(0)\rangle, \\ |\chi_1^{1,2}(T)\rangle = -iY|1\rangle = |\chi_1^{3,4}(0)\rangle, \quad (\text{B40})$$

but

$$|\chi_0^{1,2}(0)\rangle = |0\rangle \neq |\chi_0^{3,4}(T)\rangle = -i|0\rangle, \\ |\chi_1^{1,2}(0)\rangle = |1\rangle \neq |\chi_1^{3,4}(T)\rangle = i|1\rangle, \quad (\text{B41})$$

i.e., this basis is not single valued. To make it single valued, we can modify it along the third and fourth segments as

$$|\chi_0^{3,4}(t)\rangle \rightarrow |\tilde{\chi}_0^{3,4}(t)\rangle = e^{i\omega_0(t)}|\chi_0^{3,4}(t)\rangle, \\ |\chi_1^{3,4}(t)\rangle \rightarrow |\tilde{\chi}_1^{3,4}(t)\rangle = e^{i\omega_1(t)}|\chi_1^{3,4}(t)\rangle, \quad (\text{B42})$$

where

$$\omega_0(0) = 0, \quad \omega_0(T) = \frac{\pi}{2}, \quad (\text{B43})$$

$$\omega_1(0) = 0, \quad \omega_1(T) = -\frac{\pi}{2}. \quad (\text{B44})$$

Expression (B6) then becomes

$$U_{jk,jk}^\gamma = \exp\left(\int_0^T \langle \chi_j^{1,2}(t) | \frac{d}{dt} | \chi_j^{1,2}(t) \rangle dt \right. \\ \left. + \int_0^T \langle \chi_j^{3,4}(t) | \frac{d}{dt} | \chi_j^{3,4}(t) \rangle dt + (-1)^j \frac{\pi}{2} \right), \quad j = 0, 1. \quad (\text{B45})$$

But

$$\langle \chi_j^{1,2}(t) | \frac{d}{dt} | \chi_j^{1,2}(t) \rangle = -i \frac{\pi}{2T} \langle j | Y | j \rangle = 0 \quad (\text{B46})$$

and

$$\langle \chi_j^{3,4}(t) | \frac{d}{dt} | \chi_j^{3,4}(t) \rangle = i \frac{\pi}{2T} \langle j | Y X Y | j \rangle = 0. \quad (\text{B47})$$

Therefore, we obtain Eq. (B34).

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